

On the dynamics of transfer-ionization in fast ion-atomic collisions

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Abstract

We consider transfer-ionization in collisions of fast ($3.6 - 11$ MeV/u) protons, alpha-particles and lithium nuclei with helium atoms. There are just a few basic mechanisms contributing to this process which can be grouped into correlated ones, which crucially depend on the electron-electron interaction, and uncorrelated, which do not need this interaction to proceed. We show that by exploring momentum spectra of the emitted electrons the correlated and uncorrelated mechanisms can be clearly separated from each other. This exploration also enables one to get insight into subtle details of the dynamics of transfer-ionization.

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I. INTRODUCTION

Ionization and electron transfer (electron capture), which may occur in collisions between an atom and a bare nucleus, represent two of basic collision processes studied by atomic physics. In the process of ionization the atom emits an electron, which after the collision moves freely in space, while in the transfer process an electron initially bound in the atom is captured into a bound state of the moving ion. Both of these processes possess interesting physics, their study is of importance for many applications and various aspects of these processes have been attracting attention for decades.

Quite an interesting situation is encountered when a combination of transfer and ionization occurs in a single collision event. Such a process, which becomes possible if the atomic target has at least two electrons, is called transfer-ionization. During the last decade transfer-ionization in collisions of protons with helium atoms has attracted much attention [1]-[6].

There are only a few known basic mechanisms governing transfer-ionization in fast collisions. Depending on whether the electron-electron interaction (correlations) plays in them a crucial role or not, these mechanisms can be subdivided into "correlated" and "uncorrelated" ones.

The group of uncorrelated mechanisms consists of the so called independent transfer-ionization (ITI) and capture-shake-off (C-SO). In the ITI electron capture and emission occur due to the "independent" interactions of the electrons with the ionic projectile. In a theoretical description this mechanism appears starting with second order perturbation theory in the ion-atom interaction and for its realization does not need any electron-electron interaction.

According to the C-SO mechanism, a fast removal of the captured electron from the atom leads to a "sudden" change of the atomic potential in which the other electron is moving. As a result, the electron tries to adjust its state to the new potential and has a nonzero probability to become unbound [7].

The correlated mechanisms are more interesting. They include so called electron-electron Thomas (EET) mechanism and a mechanism which will be termed here as electron-electron Auger (EEA). According to the EET, transfer-ionization proceeds in two steps [8], [9]. In the first step, as a result of a binary collision with the ion, one of the electrons acquires a velocity $\sqrt{2}v$, where v is the ion velocity, moving under the angle of 45° with respect to the motion of the ion. In the second step this electron scatters on the other electron acquiring, as simple kinematics shows, a velocity equal to the projectile velocity, both in absolute magnitude and direction, that makes the capture very probable. The same kinematics also tells that the other electron in this process gets a velocity, which is perpendicular to the projectile velocity and whose absolute value is equal to v . Thus, as a result of the EET one electron is captured and the other is emitted perpendicular to the projectile motion.

The electron-electron interaction is also the (main) driving force of the EEA mechanism. The physics of the latter becomes very transparent when it is viewed in the rest frame of the projectile nucleus. The functioning of this mechanism is based on the fact that the presence of the second nucleus makes the initial configuration of atomic particles unstable with respect to a kind of Auger decay. Indeed, in the presence of this nucleus one of the electrons, which initially belongs to a bound configuration of fast moving particles constituting the atom, undergoes a radiationless transition [10] into a bound state of the ion by transferring (most of) energy excess to the another atomic electron which, as a result of this, is emitted from the atom [11], [12]. A clear signature of this mechanism is that in the rest frame of the atom the electron is emitted in the direction opposite to the projectile motion [11], [12], [13].

One has to emphasize that the mechanisms for transfer-ionization, discussed above, are in essence high-energy approximations, the validity of which improves with increasing impact energy. Therefore, the description of transfer-ionization in terms of these mechanisms becomes really

meaningful only provided the collision velocity is high enough: $v \gg v_i, v_f$, where $v_i \sim Z_t$ and $v_f \sim Z_p$ the typical velocities of the electron(s) in the initial and final bound states, respectively, and Z_t (Z_p) is the charge of the nucleus of the target (projectile). This implies that in order to get an insight into the physics of transfer-ionization by considering this process as driven by these mechanisms, even in collisions with protons the impact velocity should lie in the range $v \gtrsim 10v_0$, where v_0 is the Bohr velocity in atomic hydrogen.

Although transfer-ionization was studied in a number of papers, most of them were concerned with the total cross section. A better understanding of the physics of this process can be obtained when differential cross sections are explored. Concerning such cross sections in the case of transfer-ionization in fast collisions only the cross sections singly differential in the momentum component of the emitted electron or the target recoil ion, parallel/antiparallel to the projectile velocity, have been considered (see e.g. [6], [11]-[12]).

However, the exploration of such singly differential cross sections even in principle can hardly allow one to clearly separate the contributions of the correlated and uncorrelated mechanisms (and thus to study and understand them better). Compared to the singly differential cross sections the doubly differential cross sections, which are a function of both parallel and perpendicular to the projectile velocity components of the momentum of the emitted electron, can yield much more information about the process. Therefore, in the present paper we consider such cross sections for transfer-ionization in collisions of fast protons, alpha-particles and lithium nuclei with helium atoms. It will, in particular, be shown that the study of such doubly differential cross sections may enable one to clearly separate and identify the different mechanisms contributing to transfer-ionization and to get a better insight into the physics of this process.

One should say that all the previous experimental studies devoted to the spectra of electrons emitted in transfer-ionization were dealing with relatively low impact velocities where, as was already mentioned, the discussion of this process in terms of the four mechanisms may not yet be very meaningful. Therefore, we hope that the present article could trigger the interest of experimental physicists to the exploration of this process at higher impact velocities.

Atomic units ($\hbar = m_e = |e| = 1$) are used throughout the paper except where the otherwise stated.

II. GENERAL CONSIDERATION

In our description of transfer-ionization the correlated and uncorrelated mechanisms shall be treated separately (and added in the cross section incoherently). We begin with considering the correlated ones.

A. The EEA and EET mechanisms

The (approximate) transition amplitude for transfer-ionization can be written

$$a_{fi} = -i \int_{-\infty}^{+\infty} dt \langle \Psi_f(t) | \hat{W} | \Psi_i(t) \rangle. \quad (1)$$

Here \hat{W} is the coulomb interaction between the electrons and $\Psi_i(t)$ and $\Psi_f(t)$ are the initial and final states of the electrons.

In the nonrelativistic domain of atomic collisions the description of electron capture is covariant under a Galilean transformation and any Galilean frame can be chosen to consider this process. Assuming that the target atom is (initially) at rest in the laboratory frame, we take for the moment

the rest frame of the atom as our reference frame and choose its origin at the position of the atomic nucleus. We denote the coordinates of the electrons by \mathbf{r}_1 and \mathbf{r}_2 . The projectile-nucleus with a charge Z_p is assumed to move along a straight-line trajectory $\mathbf{R}(t) = \mathbf{b} + \mathbf{v}t$, where \mathbf{b} is the impact parameter, \mathbf{v} the collision velocity and t the time. The coordinates of the 'first' and 'second' electrons with respect to the position of the projectile are denoted by \mathbf{s}_1 and \mathbf{s}_2 , respectively ($\mathbf{s}_j = \mathbf{r}_j - \mathbf{R}(t)$; $j = 1, 2$).

We choose the initial state as

$$\Psi_i(t) = \Lambda_i \varphi_i(\mathbf{r}_1, \mathbf{r}_2) \exp(-iE_i t). \quad (2)$$

In Eq.(2) φ_i is the initial unperturbed two-electron atomic state with an energy E_i and Λ_i is a factor which accounts for the distortion of the initial atomic state by the field of the incident ion, its form shall be specified later. We approximate the state φ_i according to

$$\varphi_i(\mathbf{r}_1, \mathbf{r}_2) = A_i (\exp(-\alpha r_1 - \beta r_2) + \exp(-\alpha r_2 - \beta r_1)) \exp(\gamma r_{12}), \quad (3)$$

where A_i is the normalization factor, $r_{12} = |\mathbf{r}_1 - \mathbf{r}_2|$ is the inter-electron distance and the parameters α , β and γ can be chosen from the following sets: (i) $\alpha = \beta = 2$, $\gamma = 0$; (ii) $\alpha = \beta = 1.69$, $\gamma = 0$; (iii) $\alpha = \beta = 1.86$, $\gamma = 0.254$; (iv) $\alpha = 2.18$, $\beta = 1.18$, $\gamma = 0$; and (v) $\alpha = 2.21$, $\beta = 1.44$, $\gamma = 0.207$.

The final state is taken according to

$$\begin{aligned} \Psi_f(t) = & \Lambda_f \frac{1}{\sqrt{2}} [\chi_f(\mathbf{s}_1) \exp(i\mathbf{v} \cdot \mathbf{r}_1) \phi_{\mathbf{k}}(\mathbf{r}_2) + \chi_f(\mathbf{s}_2) \exp(i\mathbf{v} \cdot \mathbf{r}_2) \phi_{\mathbf{k}}(\mathbf{r}_1)] \\ & \times \exp(-i(\epsilon_k + \varepsilon_f)t) \exp\left(-i\frac{v^2}{2}t\right). \end{aligned} \quad (4)$$

Here, χ_f is the final (unperturbed) bound state of the electron captured by the projectile, ε_f the energy of this state (as viewed in the rest frame of the projectile) and $\exp(i\mathbf{v} \cdot \mathbf{r}_j - iv^2 t/2)$ the so called translational factor. Further, $\phi_{\mathbf{k}}$ denotes the state of the emitted electron which moves in the field of the residual atomic ion with (asymptotic) momentum \mathbf{k} and energy $\epsilon_k = k^2/2$ and Λ_f describes the distortions of the states of captured and emitted electrons by the fields of the residual atomic ion and projectile, respectively.

Now we turn to the discussion of the form of the distortion factors Λ_i and Λ_f . Let us remind the reader that in this paper we consider only collisions at high impact velocities in which one has $Z_p/v \ll 1$. Besides, as will be seen below, in the transfer-ionization process the emitted electron has a high velocity ($\sim v \gg Z_p$) with respect to the projectile. From the work on atomic ionization it is known that in such collisions the account of the distortion does not noticeably changes the result. At the same time it is also known that for electron transfer reactions the effect of the distortion in general remains very important even at $Z_p/v \ll 1$. Therefore, in our treatment we shall ignore the distortions for that electron, which is to be emitted, and account only for the distortions acting on that electron which is to be captured.

With such an assumption one can show that the transition amplitude in the momentum space,

$$S_{fi}(\mathbf{q}_\perp) = \frac{1}{2\pi} \int d^2\mathbf{b} a_{fi}(\mathbf{b}) \exp(i\mathbf{q}_\perp \cdot \mathbf{b}), \quad (5)$$

is given by

$$S_{fi}(\mathbf{q}_\perp) = S_{fi}^{\alpha,\beta}(\mathbf{q}_\perp) + S_{fi}^{\beta,\alpha}(\mathbf{q}_\perp). \quad (6)$$

Here,

$$S_{fi}^{\alpha,\beta}(\mathbf{q}_\perp) = -\frac{\sqrt{2}iA_i}{(2\pi)^3v} \int d^3\mathbf{s} \chi_f^*(\mathbf{s}) \exp(i\mathbf{q} \cdot \mathbf{s}) \Lambda_i(\mathbf{s}) \int d^3\boldsymbol{\kappa} \frac{G_{ion}(\boldsymbol{\kappa}; \beta)}{\kappa^2 + \gamma^2} \times \int d^3\mathbf{r} \Lambda_f^*(\mathbf{r}) \exp(-i(\mathbf{v} + \mathbf{q} + \boldsymbol{\kappa}) \cdot \mathbf{r}) \exp(-\alpha r) \quad (7)$$

where

$$G_{ion}(\boldsymbol{\kappa}; \beta) = \int d^3\mathbf{r} \phi_{\mathbf{k}}^*(\mathbf{r}) \exp(i\boldsymbol{\kappa} \cdot \mathbf{r}) \exp(-\beta r) \quad (8)$$

and

$$\mathbf{q} = \left(\mathbf{q}_\perp, \frac{E_i - \varepsilon_f - k^2/2 - v^2/2}{v} \right) \quad (9)$$

is the momentum transfer in the collision. Note that $S_{fi}^{\beta,\alpha}(\mathbf{q}_\perp)$ is obtained from $S_{fi}^{\alpha,\beta}(\mathbf{q}_\perp)$ by interchanging α and β in Eqs. (7)-(8).

The explicit form of the distortion factors is taken according to the continuum-distorted-state (CDW) model which has been proved quite successful in describing the total cross section for capture in a three-body collision system (one active electron + two nuclei). In this model the distortion factors read

$$\begin{aligned} \Lambda_i(\mathbf{s}) &= N(\nu_p) {}_1F_1(i\nu_p, 1, i\nu s + i\mathbf{v} \cdot \mathbf{s}) \\ \Lambda_f(\mathbf{r}) &= N^*(\nu_t) {}_1F_1(-i\nu_t, 1, -i\nu r - i\mathbf{v} \cdot \mathbf{r}), \end{aligned} \quad (10)$$

where $N(\nu) = e^{\pi\nu/2}\Gamma(1 - i\nu)$, $\nu_p = Z_p/v$, $\nu_t = Z_t/v$, and Γ and ${}_1F_1$ are the gamma and confluent hypergeometric functions, respectively (see e.g. [14]).

The inclusion of the distortion factor for the initial state in the form given by the first line of (10) means that in our treatment the electron, which is to be transferred, in its initial state moves not only in the field of the atom but also in the (coulomb) field of the projectile. Therefore, with such a factor the transition amplitude (1) describes both the EEA and EET capture channels while when this factor is set to unity, $\Lambda_i = 1$, the calculated contribution of the EEA mechanism becomes much larger but the EET mechanism simply "vanishes".

To conclude this subsection let us note that the account of the distortion for the final state turned out to be not so crucial. Indeed, in cases tested the difference between results obtained with the distortion factor Λ_f in the form given by the second line of (10) and by setting $\Lambda_f = 1$ was not substantial. Therefore, taking into account that the neglect of this distortion greatly simplifies the calculation reducing the computation time, in what follows we shall report only results obtained when we suppose that $\Lambda_f = 1$.

B. Independent transfer-ionization and capture-shake-off

Let us now very briefly consider two uncorrelated mechanisms: the independent transfer-ionization and capture-shake-off.

According to the first of them transfer-ionization proceeds in two independent steps: one electron is captured (transfer) and the other one is emitted (ionization). These transitions are driven by the interaction between the projectile and the electrons while the electrons do not need at all to interact with each other for the transitions to occur. Note that within this mechanism the projectile must interact with the target at least twice (at least one interaction per electron).

In the consideration of the present paper the capture and ionization parts of the independent transfer-ionization are regarded as occurring in the collision between a projectile-nucleus and a hydrogen-like system. The latter is described using an effective nuclear charge which was taken to be 1.69, both for capture and ionization. In the impact parameter space the amplitude for this process is a product of the single-electron transition amplitudes for capture and ionization. The latter ones are obtained using the three-body CDW (capture) and CDW-EIS (ionization) [15] models.

In capture-shake-off the "instant" removal of one of the electrons from the atom due to its capture by the fast projectile forces the other electron to react to a sudden change of the atomic potential. As a result, the second electron can be shaken off from the target and become unbound [7]. The amplitude for this channel is estimated as the product of the amplitude for single electron capture (evaluated within the three-body CDW – like in case of the ITI) and the amplitude for shake-off which is simply an overlap between the initial and final states of the "second" electron.

C. The total contribution to transfer-ionization

In our calculations we add the contributions of the correlated, the independent and capture-shake-off channels incoherently. In the context of the present paper, which is focused on the correlated capture mechanisms, such an incoherent addition does not represent a big shortcoming since at the collision parameters considered here the correlated and uncorrelated have a small overlap in the momentum space of the emitted electrons.

To conclude this section note that the validity of our approach to transfer-ionization in fast collisions has been already tested in [11]-[12] where the cross sections singly differential in the longitudinal momentum of the emitted electrons and target recoil ions were calculated for proton on helium collisions at $v = 12.6$ and 15.2 a.u. and a good agreement with available experimental data has been found [16].

III. RESULTS AND DISCUSSION

In this section we discuss the momentum spectra for electrons emitted in transfer-ionization in collisions of protons, alpha-particles and bare lithium nuclei with helium.

As was mentioned in the previous section, in our evaluation of the ITI and C-SO mechanisms we use the effective charge of 1.69 to describe the initial undistorted state of the electrons in helium. Therefore, for consistency, in our calculation of the contributions from the correlated mechanisms we use the set ii) of the parameters for the state (3) (except in figure 7, where the sets i) and v) are used).

The momentum spectra shown in figures 1-6 are given in the rest frame of the target (laboratory frame) and are represented by the doubly differential cross section

$$\frac{d^2\sigma}{dk_{lg}dk_{tr}} = k_{tr} \int_0^{2\pi} d\varphi_k \int d^2\mathbf{q}_\perp |S_{fi}(\mathbf{q}_\perp)|^2, \quad (11)$$

where $k_{lg} = \mathbf{k} \cdot \mathbf{v}/v$ and $\mathbf{k}_{tr} = \mathbf{k} - k_{lg}\mathbf{v}/v$ are the longitudinal and transverse parts, respectively, of the momentum \mathbf{k} of the emitted electron. The integration in (11) runs over the transverse part of the momentum transfer and the azimuthal angle φ_k of the emitted electron. In the range of collision parameters considered here an atomic electron is mainly captured into the ground state of the projectile. Therefore, in what follows we consider transfer-ionization only for this channel.

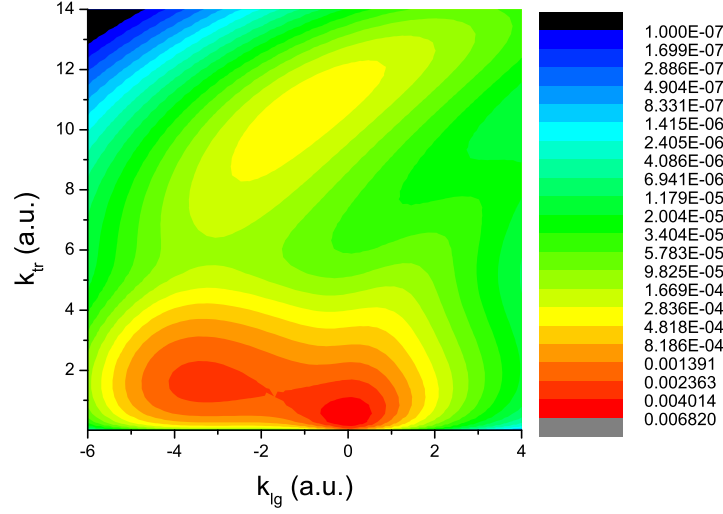


FIG. 1: Momentum spectrum (in $b/(a.u.)^2$) of electrons emitted in the reaction $3.6 \text{ MeV } p^+ + \text{He}(1s^2) \rightarrow \text{H}(1s) + \text{He}^{2+} + e^-$ collisions ($v = 12 \text{ a.u.}$).

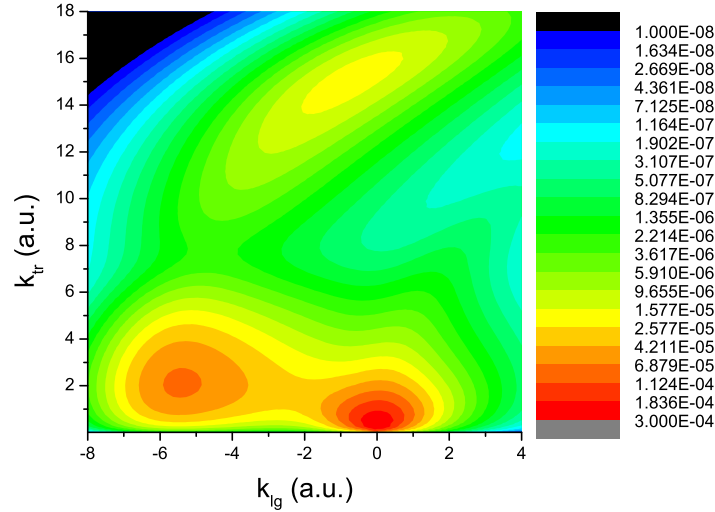


FIG. 2: Same as in figure 1 but for $6.4 \text{ MeV } p^+ + \text{He}(1s^2) \rightarrow \text{H}(1s) + \text{He}^{2+} + e^-$ collisions ($v = 16 \text{ a.u.}$).

The momentum spectra of electrons emitted in collisions with protons are displayed in figures 1, 2 and 3 for impact energies of 3.6, 6.4 and 11 MeV, respectively. These energies correspond to $v = 12$, 16 and 21 a.u. It is seen in the figures that there are three distinct maxima in the spectra.

Uncorrelated transfer-ionization

The maximum, which is located at small values of k , has its origin in the uncorrelated mechanisms: the independent transfer-ionization and capture-shake-off.

In high-velocity collisions ($v \gg Z_p, Z_t$) the cross section for single electron capture calculated within the CDW approximation scales approximately as Z_p^5/v^{11} . In our model, this is obviously also the scaling for the contribution of the capture-shake-off channel to the cross section.

Since the ionization part of the independent transfer-ionization adds the factor Z_p^2/v^2 , the cross

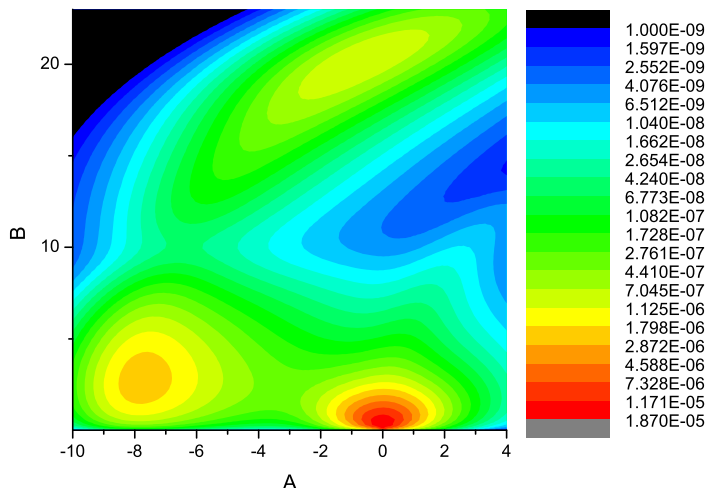


FIG. 3: Same as in figure 1 but for 11 MeV $p^+ + \text{He}(1s^2) \rightarrow \text{H}(1s) + \text{He}^{2+} + e^-$ collisions ($v = 21$ a.u.).

section for this channel is proportional to Z_p^7/v^{13} and, compared to the capture-shake-off, shows a steeper dependence both on the projectile charge and collision velocity.

According to our model, in collisions with protons (in the range of impact velocities considered) the maximum at small k is dominated by capture-shake-off, that leads to the shape of the spectrum almost symmetric with respect to $k_{lg} = 0$ [17]. The situation becomes somewhat different in collisions with alpha-particles and lithium nuclei in which the independent transfer-ionization becomes relatively more important and, as a result, the emission spectrum acquires a slight forward-backward asymmetry with more emitted electrons moving in the forward semi-sphere (see figures 4, 5 and 6).

Correlated transfer-ionization

The maximum at large (negative) k_{lg} appears due to the EEA mechanism whereas the maximum at large k_{tr} is a signature of the EET channel.

i) Let us consider the kinematics of these two correlated channels of transfer-ionization. To this end it is convenient to go first to the rest frame of the projectile-nucleus. In this frame the latter particle does not take part in the energy balance of the process (because it is heavy and is initially at rest). Therefore, the energy balance can be written as $u_e^2/2 + \Delta E \approx v^2$. Here, u_e is the velocity of the emitted electron, $\Delta E = v\Delta Q_{lg}$ is the change in energy of the nucleus of the atom with ΔQ_{lg} being the change in its longitudinal momentum, v^2 is the initial energy of the two incident electrons and we have neglected the initial and final binding energies since $v \gg Z_p$ and $v \gg Z_t$. Thus, the velocity u_e of the emitted electron in the projectile frame is approximately given by

$$u_e = v\sqrt{2(1 - \Delta Q_{lg}/v)}. \quad (12)$$

ii) If the nucleus of the atom would be just a spectator in the collision (and thus $\Delta Q_{lg} = 0$), one would obtain $u_e \approx \sqrt{2}v$. Taking into account that in the target frame the electron emitted via the EEA mechanism moves in the direction, which is opposite to the projectile velocity, the momentum spectrum of electrons produced via the EEA should then be centered in this frame around $k_{lg} \approx v - \sqrt{2}v \approx -0.4v$. Looking at the figures one sees, however, that only at the highest

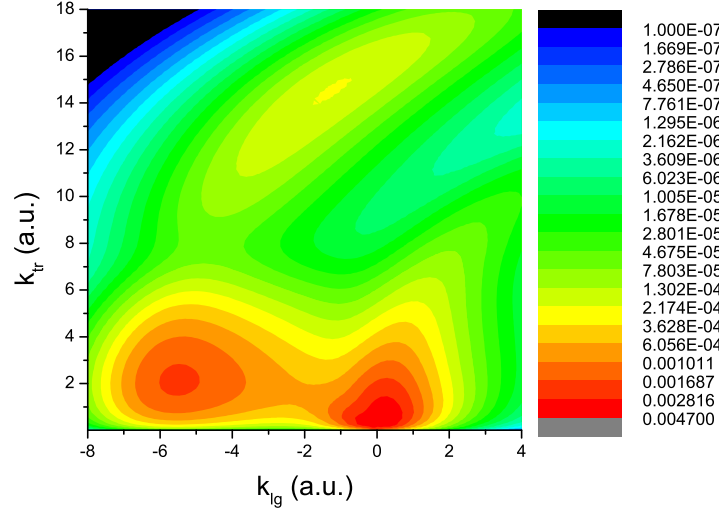


FIG. 4: Same as in figure 1 but for 6.4 MeV/u $\text{He}^{2+} + \text{He}(1s^2) \rightarrow \text{He}^+(1s) + \text{He}^{2+} + e^-$ collisions ($v = 16$ a.u.).

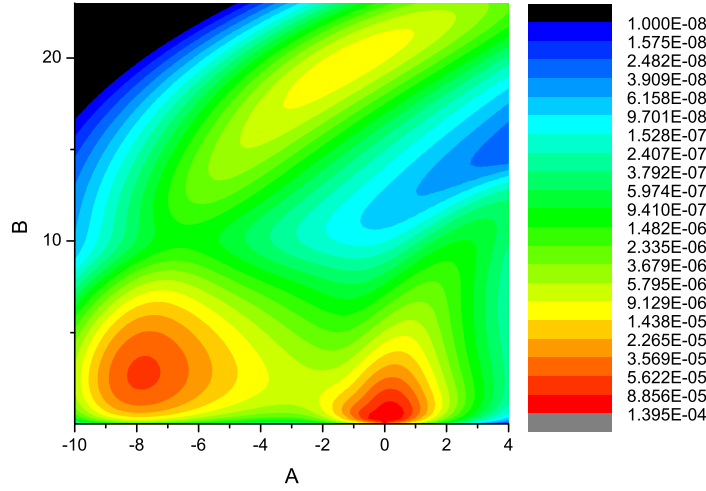


FIG. 5: Same as in figure 1 but for 11 MeV/u $\text{He}^{2+} + \text{He}(1s^2) \rightarrow \text{He}^+(1s) + \text{He}^{2+} + e^-$ collisions ($v = 21$ a.u.).

impact energy considered ($v = 21$ a.u.) the electron spectrum is really having the maximum at the longitudinal momentum rather close to $-0.4v$ while at the lower velocities ($v = 12$ and 16 a.u.) this maximum is located at a noticeable distance from the point $k_{lg} = -0.4v$. This means that only at sufficiently high impact energies the EEA becomes (almost) purely electronic mechanism without the involvement of the nucleus of the atom. At lower impact velocities the target nucleus does noticeably participate in this mechanism (see also [12]).

iii) Following the simple picture of the EET mechanism, which was mentioned in the Introduction, one would expect that in the rest frame of the target atom the EET is characterized by electrons emitted with a velocity v under the angle of 90° with respect to the motion of the projectile. Correspondingly, the velocity of the electron with respect to the projectile should be equal to $\sqrt{2}v$. The latter value indeed agrees with the simple estimate for the electron velocity u_e given above in this subsection (if we assume that $\Delta Q_{lg} = 0$).

However, according to the spectra shown in figures 1-6, in the target frame the velocity v_{EET}

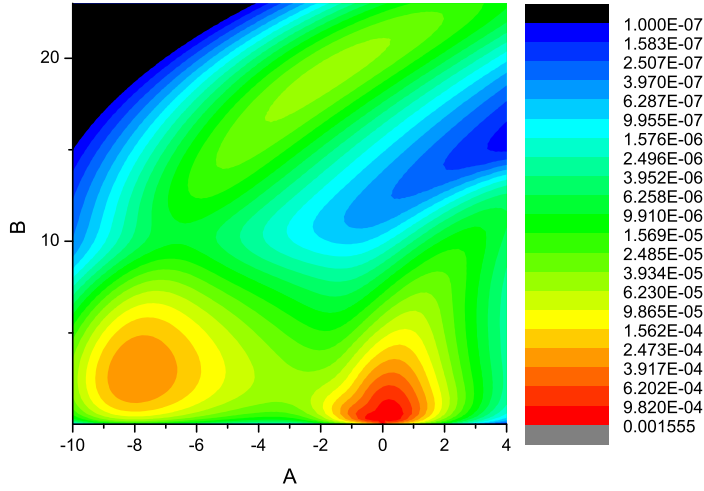


FIG. 6: Same as in figure 1 but for 11 MeV/u $\text{Li}^{3+} + {}_1^4\text{He}(1s^2) \rightarrow \text{Li}^{2+}(1s) + \text{He}^{2+} + e^-$ collisions ($v = 21$ a.u.).

of the emitted electron is on average slightly smaller than v . Besides, the angle ϑ_{EET} , which characterises the position of the “center of mass” of the EET maximum in the momentum spectrum, is somewhat larger than 90° : $\vartheta_{EET} = 90^\circ + \delta$, where $\delta > 0$. Moreover, the differences $(v - v_{EET})$ and δ increase with increasing the charge of the projectile and/or decreasing the impact velocity. The reason for this is that the simple picture does not take into account that the electron which moves together with the projectile is actually not free but bound [18]. When the binding of the captured electron increases the differences between the result and what the simple picture suggests also growth. Only in the limit $v \rightarrow \infty$ does the position of the EET maximum coincide with the prediction of this picture (see also [9]).

Assuming that at sufficiently high impact velocities the nucleus of the target atom is a spectator in the EET mechanism, one can find a simple relation between the averaged velocity v_{EET} of the emitted electron and the angle δ . Indeed, in the rest frame of the projectile the energy of this electron is approximately given by v^2 . Taking into account that the same energy can also be expressed as $(v + v_{EET} \sin \delta)^2 / 2 + (v_{EET} \cos \delta)^2 / 2 = v^2 / 2 + v_{EET}^2 / 2 + vv_{EET} \sin \delta$, one obtains $v_{EET} \approx v \left(\sqrt{1 + \sin^2 \delta} - \sin \delta \right) \approx v (1 - \delta)$.

iv) Two more observations can be drawn from figures 1-6. First, for a fixed projectile charge state the relative importance of the EET versus EEA increases with v . Second, for a fixed collision velocity v the EEA mechanism gains in relative importance when the charge of the projectile increases.

The first observation can be understood noting that the EEA and EET are basically the first and second order processes, respectively [12], [9]. As a result, the EET weakens more slowly with increasing v than the EEA. Further, the dependence of the EEA mechanism on Z_p is a bit steeper ($\sim Z_p^5$) [12] than that of the EET ($\sim Z_p^5 / (Z_p + Z_t \sqrt{2})$) [9]. This enables one to understand the second observation.

Dynamic versus stationary correlations

Both the EEA and EET mechanisms crucially rely on the coulomb interaction between the electrons. On the other hand, the electron-electron correlations in the initial and final asymptotic

states of the colliding system are also manifestation of this interaction [19].

The principal difference between them is that while the correlations in the asymptotic states are stationary in nature, the EEA and EET are based on the electron-electron interaction in its dynamic variant.

An illustrative example of the correspondence between dynamic and stationary manifestations of basically the same force is represented by the interaction between an electron and its electromagnetic field (the radiation field). A stationary situation is realized, for instance, when one considers a free (undistorted) hydrogen atom in the ground state. In this case the interaction with the radiation field has quite a weak impact on the system: it merely leads to a tiny shift of the energy of the ground state. Let us consider, however, a situation when the hydrogen atom collides with a fast ion. Now the same interaction may lead to electron transfer from the atom to the ion, which is called radiative electron capture. In this dynamic situation the interaction with the radiation field leads to a drastic change in the state of the electron.

The difference between the stationary and dynamic manifestations of the electron-electron interaction is not that dramatic. Nevertheless, it is the dynamic electron correlations (the EEA and EET), which drive the process of transfer-ionization, whereas the stationary ones providing an “environment” also influence the process by determining, for instance, the mean electron-electron distance in the initial atomic state and, thus, the magnitude of the dynamic force acting between the electrons in the transfer process [20].

These points can be seen in figure 7 where we present the contributions to the momentum spectrum due to the EEA and EET mechanisms calculated with two different approximations for the ground state of helium. In the left plot the parameters of the state (3) were taken as $\alpha = \beta = 2$ and $\gamma = 0$ which means that the electron-electron correlations in this state are completely ignored. The right plot was obtained by choosing $\alpha = 2.21$, $\beta = 1.44$ and $\gamma = 0.207$ which includes (in an approximate way) both the radial and angular correlations in the ground state of helium. It is seen that there is practically no difference in shape of these two spectra. However, their absolute intensities differ by about a factor of 2 since the electron-electron interaction in the ground state of helium increases the mean electron-nucleus distances in the atom and, thus, decreases the effective strength of the EEA and EET channels (see also [12]).

Kinematics of the uncorrelated transfer-ionization channels

To conclude our discussion in this section note that Eq. (12) is of course also valid for the uncorrelated mechanisms. In contrast to the correlated ones, however, in this case we have $u_e \approx v$ and, hence, $\Delta Q_{lg} \approx v/2$. Therefore, it is the nucleus of the atom which balances (in the projectile frame) the energy change of the captured electron, $\Delta E = v\Delta Q_{lg} \approx v^2/2$, both in the independent transfer-ionization and capture-shake-off channels.

IV. CONCLUSIONS

We have considered in some detail transfer-ionization in collisions of fast protons, alpha-particles and lithium nuclei with atomic helium. There are four basic mechanisms which are responsible for this process. Two of them (the independent transfer-ionization and capture-shake-off) are so called uncorrelated mechanisms which means that they would not disappear if the electron-electron interaction would be “switched off”. In contrast, this interaction does play a crucial role in the other two (the electron-electron and electron-electron-Thomas) mechanisms which both are governed by the dynamic electron-electron correlations.

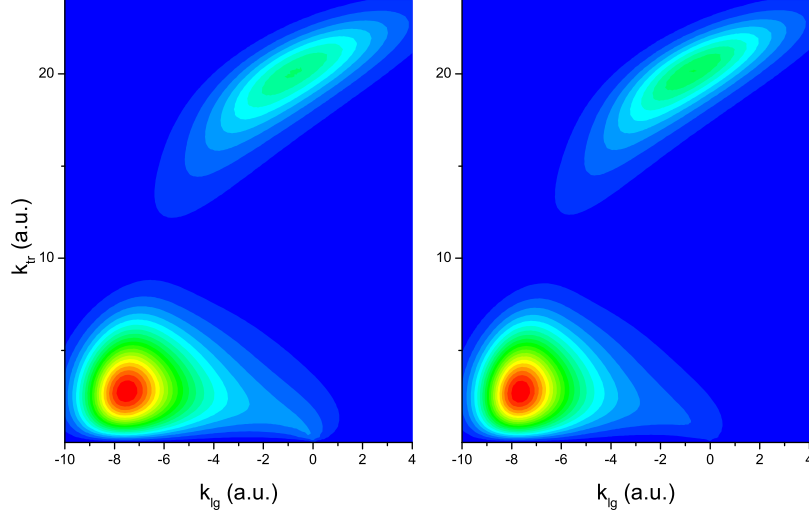


FIG. 7: The calculated contribution of the EEA and EET mechanisms to the momentum spectrum of electrons emitted in 11 MeV $p^+ + \text{He}(1s^2) \rightarrow \text{H}(1s) + \text{He}^{2+} + e^-$ collisions ($v = 21$ a.u.). The left panel: $\alpha = \beta = 2$, $\gamma = 0$. The right panel: $\alpha = 2.21$, $\beta = 1.44$ and $\gamma = 0.207$.

Our consideration shows that at sufficiently high impact velocities the contributions of the correlated and uncorrelated mechanisms can be clearly separated in the cross section doubly differential in the longitudinal and transverse components of the momentum of the emitted electron. The study of this cross section also enables one to separate the two correlated mechanisms from each other and get insight into subtle details of the dynamics of transfer-ionization.

At high impact energies $v \gg Z_t, Z_p$ the position of the center of the maximum in the momentum spectrum, caused by the EEA mechanism, tends in the target frame to $k_{lg} = -0.4v$. This means that the role of the nucleus of the atom in this mechanism weakens with increasing collision velocity and the EEA eventually becomes a truly electronic one. However, according to our model, even at impact velocities as high as 12 and 16 a.u. the helium nucleus still noticeably participates in this process.

According to the well known picture of the EET mechanism the emitted electron should have a velocity equal to the collision velocity v and fly under the angle 90° with respect to the projectile motion. Our model predicts, however, that the velocity of the emitted electron is on average smaller than v and that the electron is emitted under the angle which is larger than 90° . These two differences are interconnected and increase if the charge of the projectile increases and/or the impact velocity decreases.

An experimental exploration of the spectra of electrons emitted in transfer-ionization at high impact velocities is very desirable. At the highest velocity ($v = 21$ a.u.), considered in this article, the total cross section for transfer-ionization, according to our estimates, is of the order of 0.1, 1 and 10 mb in collisions with protons, alpha-particles and lithium nuclei. These values are of course rather small. Note, however, that already several years ago it was possible (see [6]) to measure the longitudinal momentum spectrum of the recoil target ions for transfer-ionization process with the total cross section of the order of 1 mb.

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 - [16] Moreover, even at relatively low impact velocities ($Z_p/v \simeq 0.5$) this approach can still yield reasonable results (see [13]).
 - [17] When the angular correlations are included in the ground state of helium, the calculated contribution of the C-SO channel to the emission spectrum becomes slightly asymmetric: the maximum of the spectrum is located at $p_{lg} = 0$ but more electrons are emitted in the direction opposite to the projectile motion.
 - [18] Neither this picture takes into account that the electrons are initially bound in the atom.
 - [19] A possible effect of electron correlations in the ground state of helium on transfer-ionization in fast proton-helium collisions had recently attracted quite a bit of attention [1]-[6].
 - [20] In our case the electron-electron correlations in the final state are weak due to a very high relative velocity ($\simeq \sqrt{2}v$) between the electrons.